

QCD SPECTROSCOPY AT GSI: EXOTICA AND CHARMONIA

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Abstract

In this talk I give a short summary of the basics of conventional and exotic meson spectroscopy, and consider in particular those issues in the charmonium and charmonium hybrid sectors which can be addressed by a future antiproton facility at GSI.

1 QCD and Confinement

The QCD lagrangian describes the strong interaction in terms of the couplings between the elementary pointlike quark (antiquark) and gluon constituents;

$$\mathcal{L} = \sum_q \bar{\psi}_q (i\not{D} - m_q) \psi_q - g \bar{\psi}_q (\lambda^a/2) \not{A}^a \psi_q - \frac{1}{4} G_{\mu\nu}^a G_{\mu\nu}^a . \quad (1)$$

The fundamental interactions between quarks and gluons are a QED-like $q\bar{q}g$ coupling and non-Abelian g^3 and g^4 gluon self-couplings. At large momentum scales these perturbative interactions provide an accurate description of QCD interactions, and pQCD predictions can be compared to the experimentally observed cross sections for quark and gluon jet production. At small momentum transfers however pQCD becomes inaccurate, and the dynamics of QCD is instead dominated by the nonperturbative phenomenon known as *confinement*. The gluon self-interactions lead to the formation of a “flux tube” between color sources, which gives rise to an approximately linear potential. Due to this effect only states with zero total color, “color-singlets”, can exist as physical bound states. This flux tube and the associated asymptotically linear potential are clearly evident in lattice gauge theory simulations with static sources (see Figs.1,2 [1, 2]).

2 Hadrons: Conventional and Exotica

2.1 Types of Hadrons

Hadrons are conventionally classified according to which “valence” basis state in Hilbert space is thought to dominate the hadronic state vector. This simple classification in terms of pure valence states provides a useful and surprisingly accurate description of most known resonances. Of course we should emphasize that this is an approximation of unknown accuracy with no clear justification, and may well be misleading in the description of unconventional types of hadrons.

2.2 Quark and Multiquark States

The physically allowed color-singlet states one can form from quarks and antiquarks alone are generically of the form

$$|\text{color} - \text{singlet}\rangle_{n,\bar{n}} = |q^n \bar{q}^{\bar{n}}\rangle, \quad (2)$$

where $\text{mod}(n - \bar{n}, 3) = 0$. The simplest such combinations are

$$|\text{color} - \text{singlet}\rangle_{1,1} = |q\bar{q}\rangle = |\text{quark model meson}\rangle, \quad (3)$$

$$|\text{color} - \text{singlet}\rangle_{3,0} = |qqq\rangle = |\text{quark model baryon}\rangle, \quad (4)$$

$$|\text{color} - \text{singlet}\rangle_{0,3} = |\bar{q}\bar{q}\bar{q}\rangle = |\text{quark model antibaryon}\rangle. \quad (5)$$

The complete spectrum of q and \bar{q} product basis states is formed by taking all possible quark types “flavors” for q and \bar{q} , and by allowing the states to take on all possible quark spin arrangements and orbital angular momenta, and finally by allowing excitation of the radial wavefunctions. We will discuss the detailed quantum numbers allowed to $q\bar{q}$ mesons in particular in the next section.

These three simplest color-singlet states, $|q\bar{q}\rangle$, $|qqq\rangle$ and $|\bar{q}\bar{q}\bar{q}\rangle$ are special in that they are *irreducible*, in other words they cannot be partitioned into separate color-singlet substates. The “higher” Fock space color singlets in contrast are *reducible*, and need not be realized in nature as isolated resonances. Two examples of such higher Fock space states are

$$|\text{color} - \text{singlet}\rangle_{2,2} = |q^2 \bar{q}^2\rangle = |\text{quark model baryonium}\rangle \quad (6)$$

and

$$|\text{color} - \text{singlet}\rangle_{6,0} = |qqqqqq\rangle = |\text{quark model dibaryon}\rangle. \quad (7)$$

The multiquark combination

$$|\text{color} - \text{singlet}\rangle_{4,1} = |qqqq\bar{q}\rangle \quad (8)$$

is also an allowed color-singlet basis state, but has received rather less theoretical attention.

Since these hypothetical baryonia and dibaryons have overlap with scattering states of two separate $|q\bar{q}\rangle$ mesons and two separate $|qqq\rangle$ baryons respectively, they can “decay” without interaction. It is expected therefore that they either have extremely broad widths from “*fall-apart*” into these final states, or may not be realized in nature as resonances at all. This fall-apart problem would be circumvented by a multiquark state with a mass below all strong decay thresholds, which would therefore be strongly stable. Possibilities for strongly stable multiquark states include the $u^2d^2s^2$ “H dibaryon” [3] and “heavy-light” $Q^2\bar{q}^2$ clusters, with $Q = c$ or b [4].

Alternatively, one may find quasinuclear bound states of largely undistorted hadrons that formally lie in multiquark sectors of Hilbert space, such as nuclei, hypernuclei, and perhaps $K\bar{K}$ bound states [5].

2.3 Quarkonium and $q\bar{q}$ Quantum Numbers

Most known mesons are reasonably well described as $q\bar{q}$ (quark-antiquark) bound states. Since quarks have $S = 1/2$, the $q\bar{q}$ pair can have total spin $S_{q\bar{q}} = 1/2 \otimes 1/2 = 1 \oplus 0$. The $q\bar{q}$ orbital angular momentum $L_{q\bar{q}}$ can take on any integer value; combining these L and S $q\bar{q}$ angular momenta gives the allowed total angular momentum $J_{q\bar{q}}$. The allowed values are $J = L$ (for $S = 0$) and $J = L + 1, L, L - 1$ (for $S = 1$). Meson quark model assignments may be specified using spectroscopic notation, $^{2S+1}L_J$. As examples, the π is a 1S_0 state, the J/ψ is 3S_1 and the $L = 1, S = 1, J = 2$ $f_2(1270)$ is a 3P_2 $q\bar{q}$ quark model state. Radial excitation may be indicated using a prefactor, thus the first radially-excited $1^{--} c\bar{c}$, known as the $\psi(3686)$, is a 2^3S_1 state.

Spatial parity P and charge-conjugation parity C are exact symmetries of the QCD lagrangian, and as such are conserved in strong decays. In $q\bar{q}$ states these quantum numbers are

$$P = (-1)^L \quad (9)$$

and

$$C = (-1)^{L+S} . \quad (10)$$

A state’s J^{PC} quantum numbers follow directly from these relations; for example the 1S_0 π^0 has $J^{PC} = 0^{-+}$, the 3S_1 J/ψ has $J^{PC} = 1^{--}$, and the 3P_2

$f_2(1270)$ has $J^{PC} = 2^{++}$. As we shall see, sensitive tests of the nature of interquark forces are possible given accurate experimental information on the spectrum of $q\bar{q}$ states in heavy quark systems; $c\bar{c}$ is an especially clear case.

If we complete a table of all possible J^{PC} quantum numbers allowed to $q\bar{q}$ states, we find that certain combinations do not arise. These “ J^{PC} -exotics” are $0^{--}; 0^{+-}, 1^{-+}, 2^{+-}, 3^{-+} \dots$. Since relatively low-mass J^{PC} -exotics are expected in the hybrid meson spectrum, and if discovered would certainly constitute proof of states beyond the naive $q\bar{q}$ quark model, the search for such states is widely regarded as the most exciting topic in QCD spectroscopy.

2.4 Glueball and Hybrid States

Since this subject was covered extensively by other Hirschegg Workshop speakers in the context of light (u, d, s, g) spectroscopy, I will be very brief here and proceed to the $c\bar{c}$ and $c\bar{c}$ -hybrid states.

One may also form physically allowed (color-singlet) basis states from pure gluons and from quark, antiquark and gluon product states. Hadrons which have such configurations as valence basis states are known collectively as “gluonic hadrons”. Color-singlet gluon states (if we neglect quarks) form idealized, unmixed “glueball” resonances. The spectrum of these states has been studied extensively in LGT (see for example [6]), and an impressively detailed theoretical spectrum has been determined. Unfortunately for experimentalists, LGT predicts only one glueball below 2 GeV, a scalar at 1.6 GeV, and no J^{PC} -exotics are expected below ca. 4 GeV. This scalar has been identified with the $f_0(1500)$ seen in $p\bar{p}$ annihilation [7], although there are outstanding problems with clear violation of flavor symmetry in the strong decay branching fractions of this state; naively one would expect a glueball to couple symmetrically to all quark flavors. This flavor-symmetry violation may indicate that glueball-quarkonium mixing is an important effect [7]. This mixing could be important both in decays and in mass shifts of the observed resonances relative to LGT predictions; since this is in effect a systematic error for LGT, it will be very important to quantify.

The relatively narrow width of the $f_0(1500)$ glueball candidate is very encouraging for higher-mass glueball searches; LGT predicts these to be a 0^{-+} and a 2^{++} , just below 2.5 GeV. Of course identification of these states will require clarification of the higher quarkonia expected in the same mass region. A future $p\bar{p}$ machine could make a very useful contribution through an accurate determination of the branching fractions of the $f_0(1500)$ and other light glueball and hybrid candidates.

Hybrids are the most experimentally attractive of the anticipated non- $q\bar{q}$ resonances, because their valence $|q\bar{q}g\rangle$ basis states span complete flavor nonets (hence hybrids have a much richer spectrum of states than glueballs) and the lowest-lying hybrid multiplet is expected to contain exotic quantum numbers. All J^{PC} combinations can be formed from $|q\bar{q}g\rangle$ states, so any J^{PC} -exotic might *a priori* be a hybrid meson candidate. LGT has recently contributed several estimates of light exotic hybrid meson masses [8, 9], and at present it appears that the combination $J^{PC} = 1^{-+}$ is the lightest exotic, with a mass of $M(1^{-+}) \approx 2.0$ GeV in the (u, d) $q\bar{q}$ flavor sector. This is consistent with estimates using the flux-tube model [10], and rather heavier than bag model results, which favored $M(1^{-+}) \sim 1.5$ GeV.

We can expect to identify the J^{PC} -exotic hybrids rather easily, provided that we study their favored decay modes. The expectation of both flux tube [10] and constituent gluon [11] models is that light hybrids should decay preferentially to pairs of $q\bar{q}$ mesons in which one has a unit of orbital excitation, the so-called “S+P” modes. For the lightest $I = 1$, $J^{PC} = 1^{-+}$ hybrid these modes are πf_1 and πb_1 , which are rather difficult to reconstruct and so had not been investigated carefully before the recent interest in hybrid mesons.

As a caution we note that the only two light exotic hybrid candidates, the $\pi_1(1400)$ and $\pi_1(1600)$, both have these $I = 1$, $J^{PC} = 1^{-+}$ quantum numbers, but lie ca. 500 MeV below the LGT and flux-tube mass predictions and apparently decay strongly to the S+S modes $\eta\pi$ and $\rho\pi$, which are forbidden to hybrids in the flux-tube decay model. Evidently these theoretical expectations for hybrids should not be regarded as more than tentative guidelines at present. We can of course expect a systematic improvement in LGT predictions as algorithms and computer performance improve.

3 Charmonium and $c\bar{c}$ Hybrids

3.1 Charmonium, Theory and Experiment

There are 11 known $c\bar{c}$ resonances; the spectrum is shown in Fig.3. (A possible 12th $c\bar{c}$ state, a candidate for the anticipated narrow 3D_2 , was reported in $J/\psi\pi^+\pi^-$ at 3.836 GeV by the E705 Collaboration [12]. This effect has a rather low 2.8σ statistical significance, and needs confirmation.) There is a predominance of $J^{PC} = 1^{--}$ vector states simply because most of these resonances were found at e^+e^- machines, which form only $J^{PC} = 1^{--}$ states in s-channel. The remaining states in the 1P multiplet and the 1S_0 spin-singlet $\eta_c(2980)$ were found in radiative transitions from the 1S and 2S vectors, with

the single exception of the $h_c(3526)$; this $J^{PC} = 1^{-+}$ state has been observed only in $p\bar{p}$ annihilation.

Since charmonium is only quasirelativistic, one can expect that the level of configuration mixing is much reduced relative to light mesons (the gluon emission amplitude is $\propto (v_q/c_q)$), so that one may be able to clearly identify the separate effects of confinement and gluon exchange ($|c\bar{c}\rangle \leftrightarrow |c\bar{c}g\rangle$ mixing) in the $c\bar{c}$ spectrum. Since spin-dependent forces do not appear in the effective interquark interaction until $O(v^2/c^2)$ in the quark momenta, we might expect that a naive zeroth-order static potential that incorporated the OGE color-Coulomb potential and linear confinement might give a reasonable approximation to the observed $c\bar{c}$ spectrum [13]. We can test this by assuming the potential

$$V_{c\bar{c}}(r) = -\frac{4}{3} \frac{\alpha_s}{r} + br \quad (11)$$

and solving the nonrelativistic Schrödinger equation for bound states in this potential. Inspection of the experimental $c\bar{c}$ spectrum (Fig.3) suggests 1S, 1P and 2S multiplets with spin-averaged masses of ca. 3.07 GeV, 3.52 GeV and 3.67 GeV, respectively. If we use these as input to fix the three potential model parameters α_s, b and the charm quark mass m_c , we find $\alpha_s = 0.510$, $b = 0.152 \text{ GeV}^2$ and $m_c = 1.450 \text{ GeV}$. Most of the remaining $c\bar{c}$ levels predicted to lie below 4.6 GeV (1S..4S, 1P..3P, 1D..3D, 1F, 2F, 1G, 2G and 1H) are shown in Fig.4, together with the experimental spectrum. With a multiplicity of 2 for S-states ($^1S_0, ^3S_1$) and 4 for higher- L states, this model predicts 52 independent $c\bar{c}$ states below 4.6 GeV. The proximity of the experimental masses to the predicted radial and orbital levels in Fig.4 confirms that the simple description of $c\bar{c}$ states as nonrelativistic fermions in a Coulomb-plus-linear potential is a reasonable first approximation.

The level splittings within an orbitally-excited multiplet such as 1P provide more sensitive tests of the nature of interquark forces. Assuming that the short-ranged force is due to one-gluon exchange (OGE), we expect the spin-dependent forces to be reasonably well described by the Breit-Fermi Hamiltonian, which follows from an $O(v^2/c^2)$ expansion of the OGE T-matrix. This Breit-Fermi interaction, which is familiar from atomic physics, has spin-spin, spin-orbit and tensor terms, and for equal-mass quarks and antiquarks ($m_q = m_{\bar{q}} = m$) is explicitly

$$H_{Breit-Fermi}^{OGE} = \frac{32\pi\alpha_s}{9m^2} \vec{S}_q \cdot \vec{S}_{\bar{q}} \delta(\vec{x}) + \frac{2\alpha_s}{m^2 r^3} \vec{L}_{q\bar{q}} \cdot \vec{S}_{q\bar{q}} + \frac{4\pi\alpha_s}{m^2 r^3} (\vec{S}_q \cdot \hat{r} \vec{S}_{\bar{q}} \cdot \hat{r} - \frac{1}{3} \vec{S}_q \cdot \vec{S}_{\bar{q}}) \quad (12)$$

There are also spin-dependent terms which change the relative positions of multiplets by $O(v^2/c^2)$ but do not separate levels within a multiplet.

Several very characteristic features of this OGE interaction are immediately evident. First, since the spin-spin interaction is a contact term, it has no effect on orbitally-excited states. Thus the spin-singlet state ($S = 0, J = L$) is predicted to be degenerate with the multiplicity-weighted “center-of-gravity” of the spin-triplet states ($S = 1, J = L + 1, L, L - 1$). (The spin-orbit and tensor mass shifts give zero when weighted by multiplicity.) Thus in the $1P$ multiplet we predict that the 1P_1 h_c spin singlet should have a mass of

$$M(^1P_1) \Big|_{thy.(OGE)} = \frac{5}{9} M(^3P_2) + \frac{3}{9} M(^3P_1) + \frac{1}{9} M(^3P_0) = 3525.27(0.12) \text{ MeV}. \quad (13)$$

This relation is very well satisfied by the experimental candidate $h_c(3526)$ reported by the Fermilab collaboration E760/835 in $p\bar{p}$ annihilation [14]; it has a mass of

$$M(^1P_1) \Big|_{expt.} = 3526.14(0.24) \text{ MeV}. \quad (14)$$

The agreement is not expected to be exact because this is an $O(v^2/c^2)$ derivation, and makes additional approximations such as assuming pure $c\bar{c}$ states and *only* OGE at small r .

This result is often cited as a sensitive test of the Lorentz nature of the confining interaction. *A priori* one might have assumed that the confining interaction couples to the color charge density $\psi^\dagger\psi = \bar{\psi}\gamma_0\psi$, so that the complete quark-antiquark interaction is of the same form as the QED Coulomb interaction, $\gamma_0 \otimes \gamma_0$. This was assumed in the original Cornell model, and is still advocated by Swanson and Szczepaniak [15]. Alternatively the confining interaction might couple to the Lorentz scalar density $\bar{\psi}\psi$, so that the complete interaction transforms as $I \otimes I$. These two possibilities may be distinguished by the $O(v^2/c^2)$ spin-dependent Hamiltonian. The general result for the spin-spin $c\bar{c}$ interaction due to a $\gamma_0 \otimes \gamma_0$ potential $V(r)$ is

$$H_{vector}^{spin-spin} = +\frac{2}{3m_c^2} \nabla^2 V(r) \vec{S}_q \cdot \vec{S}_{\bar{q}}. \quad (15)$$

With a *vector* linear confining interaction $V(r) = b_v r$ this becomes

$$H_{vector \text{ conf.}}^{spin-spin} = +\frac{4b_v}{3m_c^2 r} \vec{S}_q \cdot \vec{S}_{\bar{q}}. \quad (16)$$

This vector confinement would displace the 1P_1 h_c state upwards in mass from the 3P_J c.o.g. by $(4b_v/3m_c^2) \langle 1P | r^{-1} | 1P \rangle \approx 30 \text{ MeV}$; since these energies

are actually equal to within about 1 MeV, this is a very strong argument in favor of scalar over vector confinement. Since the $h_c(3526)$ state is not very well established experimentally, and has been seen only in $p\bar{p}$ annihilation, confirmation of this 1P_1 state and a precise mass determination will be a very important exercise for GSI.

In addition to this OGE interaction there is an inverted spin-orbit interaction due to the linear scalar confining potential, which is given by

$$H_{\text{scalar conf.}}^{\text{spin-orbit}} = -\frac{\pi b}{2m_c^2 r} \vec{L}_{q\bar{q}} \cdot \vec{S}_{q\bar{q}} . \quad (17)$$

The effect of incorporating these additional spin-dependent terms as first order perturbations is shown in Fig.5. Note that this is *not* a fit to the observed multiplet splittings, rather these splittings follow from the α_s, b and m_c that fit the 1S, 1P and 2S multiplet centers of gravity (Fig.4). As a qualitative description of the relative positions and scale of splittings within the multiplet this model is evidently quite successful. The relative splitting $(^3P_2 - ^3P_1)/(^3P_2 - ^3P_0)$ is an especially interesting quantity, since the wave-function uncertainties approximately cancel and one can see the effects of the negative (scalar confinement) spin-orbit and (small) OGE tensor terms. The observed ratio is rather close to theoretical expectations from OGE and linear scalar confinement. Similar tests in the 2P and especially 1D multiplets would be quite interesting at GSI, since the negative scalar spin-orbit term is longer ranged than the OGE spin-orbit, so we expect to see considerable narrowing of the multiplet splittings with increasing orbital and radial excitations. Complete inversion is predicted with increasing L , but other effects such as configuration mixing and coupling to open charm states may mask this interesting effect. Accurate mass determinations of many conventional $c\bar{c}$ states above open charm threshold will be very useful for theorists trying to quantify the various mass shifts.

Identification of the $c\bar{c}$ spectrum above the open-charm threshold at 3.73 GeV will be interesting for tests of decay models, spectroscopy models, and also because these states are a “background” which might otherwise be confused with charmonium hybrids, charm meson molecules, or other unusual states. First, the two 1D states 1D_2 (2^{-+}) and 3D_2 (2^{--}) are especially interesting because they cannot decay to DD , and hence should be relatively narrow. (DD is of course an abbreviation for $D\bar{D}$ in this context.) Studies of the relative branching fractions of other higher-mass $c\bar{c}$ states to the presumably dominant open-charm modes ($DD, D^*D, D^*D^*, D_sD, D_s^*D, \dots$) will be an extremely interesting contribution to our understanding of strong QCD physics.

Theorists usually treat open-flavor strong decays using the “ 3P_0 decay model” or one of its variants such as the flux-tube model. This type of decay model, which predates QCD, describes open-flavor decays as due to $q\bar{q}$ pair creation with vacuum (3P_0) quantum numbers. Just why this model works is unclear, and the evidence supporting it is rather meagre; the classic tests are the D/S amplitude ratios in the two decays $b_1 \rightarrow \omega\pi$ and $a_1 \rightarrow \rho\pi$ [16]. Since so much of hadron spectroscopy makes use of this decay model (the weak πN solution of the “missing baryon” problem and the S+P hybrid signature are two examples), it is very important to test it using a wide range of resonance quantum numbers and final states. The higher-mass $c\bar{c}$ states will be very useful in this regard. Calculations of the open-charm branching fractions of higher $c\bar{c}$ states have previously been reported using the 3P_0 model [17], and these show interesting dependence on the nodal structure of the radial $c\bar{c}$ wavefunctions. An experimental determination of the strong decay amplitudes of the accessible higher-mass $c\bar{c}$ states would allow an extremely interesting test of this widely used but inadequately tested strong decay model.

The little that is known about strong decays of the higher-mass $c\bar{c}$ states already includes a famous puzzle; the $\psi(4040)$, which has a mass consistent with a 3S state (Fig.4), purportedly has relative branching fractions of $D^*D^* \gg D^*D \gg DD$, despite the fact that the D^*D^* mode has essentially no phase space! This led to speculations that the $\psi(4040)$ might be a D^*D^* molecule [18], or that it may be the expected 3 3S_1 $c\bar{c}$ state, but with nearby decay amplitude zeros that lead to these anomalous branching fractions [17]. Since we hope to use branching fractions to characterize states, an accurate test of these and other strong branching fractions would clearly be a first priority at a new $p\bar{p}$ charmonium facility.

3.2 Charmonium Hybrids

The charmonium system is an excellent laboratory for the study of nonperturbative QCD effects such as confinement and gluonic excitations. It has the advantage of being quasirelativistic; the adiabatic $c\bar{c}$ potential is clearly evident in the spectrum of states, but the $O(v/c) |c\bar{c}g\rangle$ gluonic configuration mixing is sufficiently large to be accurately determined and compared with model predictions, for example in the $O(v^2/c^2)$ spin-dependent multiplet splittings of the 1P states. The simplicity of the known $c\bar{c}$ spectrum suggests that it may be straightforward to identify relatively unmixed charmonium hybrids as “extra” charmonium states through a more complete determination of the experimental spectrum. Of course the identification of complete hy-

brid multiplets, especially J^{PC} exotics, would be a crucial contribution to our understanding of the dynamics of gluonic excitations. In the charmonium system these states may be narrow enough to make this a feasible experimental program.

Recent theoretical advances have simplified the problem of searching for hybrid charmonium considerably. Previous model estimates of the mass of the lowest hybrid charmonium multiplet varied over a rather wide range, ca. 4.0-4.5 GeV. (For a review of this earlier work see Ref.[19].) With the development of lattice NRQCD we now have lattice results for the masses of exotic $c\bar{c}$ - and $b\bar{b}$ -hybrids that report very small statistical errors of ca. 10 MeV. (The systematic uncertainties are not yet known but might be ca. 50 MeV, and will be estimated in subsequent work.) As one example, the CP-PACS collaboration [9] quote masses for the 1^{-+} (expected to be the lightest exotic) $c\bar{c}$ and $b\bar{b}$ hybrid states of

$$M_{c\bar{c}H}(1^{-+}) = M_{c\bar{c}}(1S) + 1.323(13) \text{ GeV} \approx 4.39 \text{ GeV} \quad (18)$$

and

$$M_{b\bar{b}H}(1^{-+}) = M_{b\bar{b}}(1S) + 1.542(8) \text{ GeV} \approx 10.99 \text{ GeV} . \quad (19)$$

(I assume multiplicity-weighted 1S masses of 3.07 GeV for $c\bar{c}$ and 9.45 GeV for $b\bar{b}$.) Thus we have a presumably accurate lattice estimate of the mass of the lightest $c\bar{c}$ hybrid multiplet, ≈ 4.4 GeV. The precise mass of the lightest hybrid multiplet has previously been of great interest because of the flux-tube model prediction that S+P modes should be strongly favored for hybrids; if $c\bar{c}$ -hybrids were below the S+P threshold of $M(D) + M(D_j^*) \approx 4.25$ GeV, one might have anticipated relatively narrow states. With the NRQCD lattice results it now appears that S+P modes are indeed open, so $c\bar{c}$ -hybrids need not be anomalously narrow. In any case the observation of important $\pi\eta$ and $\pi\rho$ modes for the hybrid candidates $\pi_1(1400)$ and $\pi_1(1600)$ suggests that experiment may not support this selection rule as an especially rigorous one; the simple S+S modes DD , DD^* and D^*D^* may well be the dominant hybrid modes. It will be very important experimentally to search all allowed quasi-two-body open charm modes for these states. Just as with the conventional $c\bar{c}$ states, much is speculated but very little is known about open-flavor strong branching fractions of hybrids.

4 Summary and Conclusions

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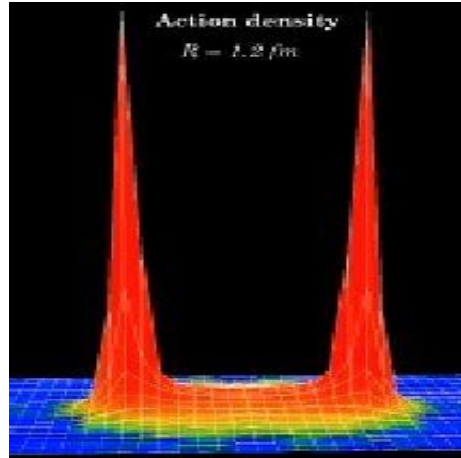


Fig. 1: The gluonic flux tube between static quark sources (Bali et al. [1]).

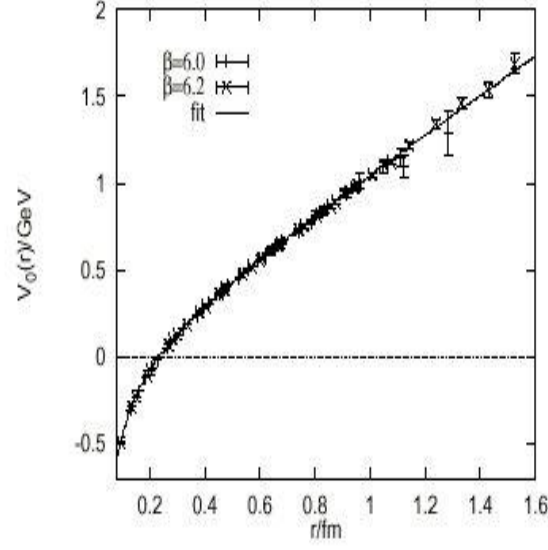


Fig. 2: An LGT determination of the interquark potential (Bali et al. [2]).

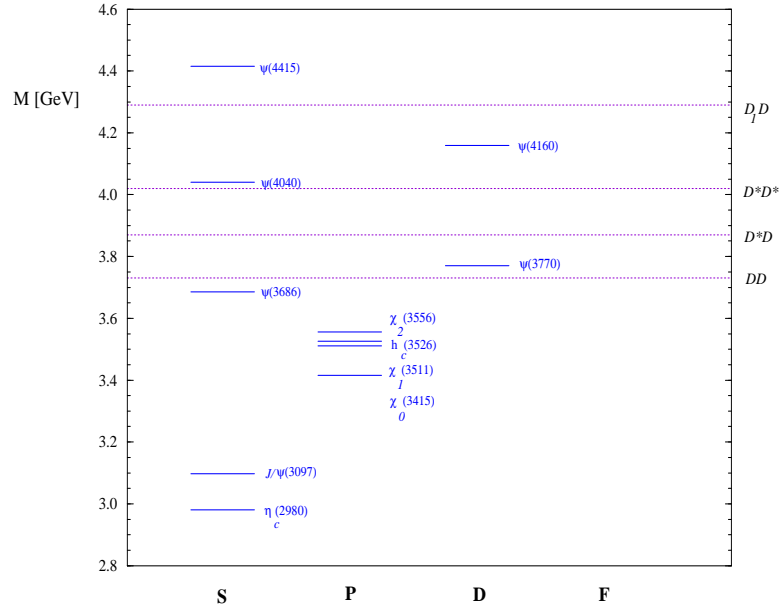


Fig. 3: The 11 known $c\bar{c}$ states. All have $J^{PC} = 1^{--}$ except the 1P multiplet near 3.5 GeV and the $J^{PC} = 0^{-+}$ $\eta_c(2980)$. Some open charm thresholds are also shown.

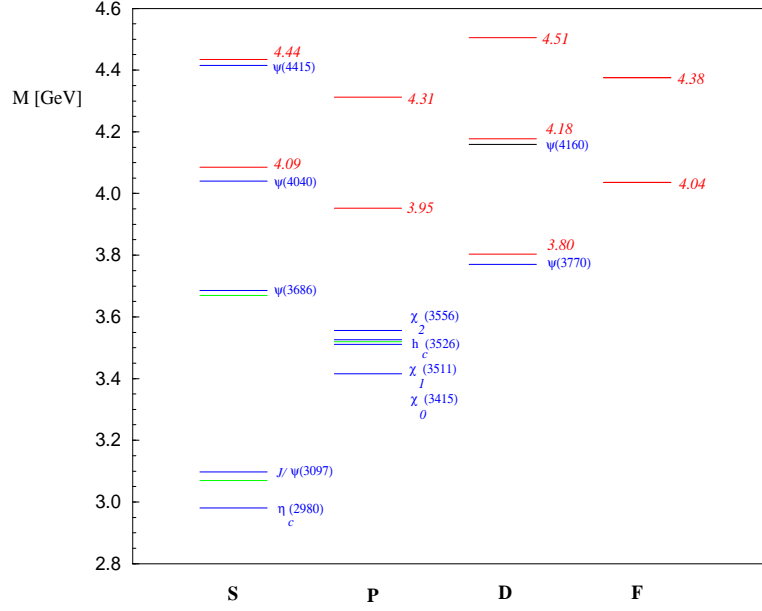


Fig. 4: A comparison of theory (red, 52 states) and experiment (blue, 11 states) for $c\bar{c}$ levels below 4.6 GeV. Mean 1S, 1P and 2S (*estm.*) levels (green) were input. Predicted levels not shown are 1G(4.24), 2G(4.56) and 1H(4.43).

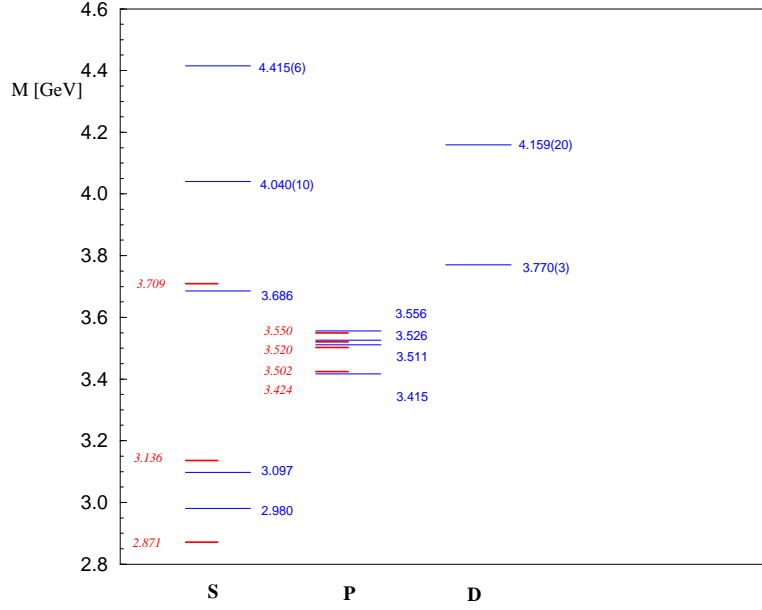


Fig. 5: Spin-dependent splittings from OGE and linear scalar confinement. Theory (red) is compared to experiment (blue). This is *not* a new fit; the parameters are those of Fig.4.